HOW LONG CAN TINY H I CLOUDS SURVIVE?

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ABSTRACT

We estimate the evaporation timescale for spherical H I clouds consisting of the cold neutral medium surrounded by the warm neutral medium. We focus on clouds smaller than 1pc, which corresponds to *tiny H I clouds* recently discovered by Braun & Kanekar and Stanimirović & Heiles. By performing one-dimensional spherically symmetric numerical simulations of the two-phase interstellar medium (ISM), we derive the timescales as a function of the cloud size and of pressure of the ambient warm medium. We find that the evaporation timescale of the clouds of 0.01 pc is about 1Myr with standard ISM pressure, $p/k_B \sim 10^{3.5}$ K cm^{-3}, and for clouds larger than about 0.1 pc it depends strongly on the pressure. In high pressure cases, there exists a critical radius for clouds growing as a function of pressure, but the minimum critical size is ~ 0.03 pc for a standard environment. If tiny H I clouds exist ubiquitously, our analysis suggests two implications: tiny H I clouds are formed continuously with the timescale of 1Myr, or the ambient pressure around the clouds is much higher than the standard ISM pressure. We also find that the results agree well with those obtained by assuming quasi-steady state evolution. The cloud-size dependence of the timescale is well explained by an analytic approximate formula derived by Nagashima, Koyama & Inutsuka. We also compare it with the evaporation rate given by McKee & Cowie.

Subject headings: hydrodynamics – ISM: clouds – ISM: kinematics and dynamics

1. INTRODUCTION

Recently discovered "tiny H I clouds" have very small sizes, $\sim 10^{-2}$ pc, and small H I column density, $\sim 10^{18}$ cm^{-2} (Braun & Kanekar 2005; Stanimirović & Heiles 2005). This new population of clouds is providing a challenge to our understanding of the evolutionary cycle of the interstellar medium (ISM). The origin of the clouds is still under debate, but it is presumably the thermal instability (Field 1965; Balbus 1986) in turbulent gas (e.g. Vázquez-Semadeni, Passot & Pouquet 1995; Koyama & Inutsuka 2002, 2004; Kritsuk & Norman 2002a,b; Hennebelle & Audit 2005) and/or the fragmentation of cold clouds crushed by interstellar shocks (Nakamura et al. 2005). Apart from the formation process, we can extract many information about the surrounding ISM via investigation of the evolution and statistics of the clouds. In this *Letter*, we thus study the evaporation rate and timescale of the tiny H I clouds, which are key quantities characterizing the fate of the clouds.

One of the simplest model is the clouds that consist of the cold neutral medium (CNM) with temperature $T \sim 10^{1-2}$ K, surrounded by the warm neutral medium (WNM) with $T \sim 10^4$ K, under pressure equilibrium. The two phases are thermally stable balancing the heating rate due to, e.g., photo-electric heating of dust grains, with the cooling rate due to, e.g., line cooling by fine-structure transition of C II atoms. This is based on a widely accepted two-phase description of the ISM (Field, Goldsmith & Habing 1969; Wolfire et al. 2003). The evolution of clouds is, therefore, described as the motion of the interface, or *front*, between the CNM and WNM, driven by the thermal conduction. Based on this picture, Zel'dovich & Pikel'ner (1969) and Penston & Brown (1970)

considered isobaric, steady motion of fronts in plane-parallel geometry. They clarified that the motion is determined by pressure and that there is a *saturation* pressure for which a static front can exist. Their work has been substantially extended by many authors (e.g. Elphick, Regev & Spiegel 1991; Elphick, Regev & Shaviv 1992; Ferrara & Shchekinov 1993; Hennebelle & Pérault 1999).

Compared to the analysis of the frontal motion in plane-parallel geometry, the evolution of spherical clouds has been much less analyzed. Graham & Langer (1973) extended the work by Zel'dovich & Pikel'ner (1969) and Penston & Brown (1970) to numerically compute isobaric flows in three dimensional spherical geometry. Cowie & McKee (1977) and McKee & Cowie (1977, hereafter MC77) considered a model of spherical clouds surrounded by the hot ionized medium and derived an analytic formula of the evaporation rate of clouds. MC77 also considered cold clouds surrounded by the WNM, which can be directly compared with the work by Graham & Langer (1973). In our previous paper, Nagashima, Koyama & Inutsuka (2005), we considered the growth of spherical clouds and showed a new way to obtaining approximate analytic formula of the evaporation and condensation rate. However, because we adopted a simple polynomial form of the cooling function proposed by Elphick, Regev & Shaviv (1992) for simplicity, our understandings were only qualitative. In this *Letter*, we shall show the evaporation rate for a realistic cooling function by using a full numerical simulation and compare it with the approximate analytic formula and with the result of MC77. We would like to stress that our analysis can be applied to any spherical cold clouds surrounded by warm gas under pressure equilibrium irrespective of their origin. Thus the results shown here are quite general.

Our aim is thus to estimate the evaporation timescale of cold H I clouds as a function of the cloud size and the pressure of the ambient WNM. This paper is outlined as follows. In §2 we compute the evaporation rate of cold clouds by using a full numerical simulation and compare it with MC77's. In §3 we

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show the evaporation timescale. In §4 we discuss the dependence of the timescale on the size and pressure. We derive a new description of the evaporation rate as an improvement upon MC77. §5 we provide conclusions.

2. EVAPORATION RATES

Below we assume that clouds are spherically symmetric for simplicity, so that we compute only radial structure of clouds. Thus, the basic fluid equations are written as

$$\frac{\partial \rho}{\partial t} + \frac{1}{r^2}\frac{\partial}{\partial r}r^2\rho v = 0, \tag{1}$$

$$\frac{\partial v}{\partial t} + v\frac{\partial v}{\partial r} = -\frac{1}{\rho}\frac{\partial p}{\partial r}, \tag{2}$$

$$\frac{\partial \rho e}{\partial t} + \frac{1}{r^2}\frac{\partial r^2 \rho e v}{\partial r} + \frac{p}{r^2}\frac{\partial r^2 v}{\partial r} = -\rho\mathcal{L} + \frac{1}{r^2}\frac{\partial}{\partial r}r^2\kappa(T)\frac{\partial T}{\partial r} \tag{3}$$

and the equation of state, $p = \rho k_B T/\mu$, where $\kappa(T) = 2.5 \times 10^3\sqrt{T}$ erg^{-1} cm^{-1} K^{-1} s^{-1} is the conductivity for neutral gas (Parker 1953), k_B the Boltzmann constant, μ the mean molecular mass so that $\rho = \mu n$, and $\mathcal{L}$ is the heat-loss function defined as $\rho\mathcal{L} \equiv n^2\Lambda - n\Gamma$, where Λ and Γ are the cooling and heating rates, respectively. We adopt a simple analytic fitting function given by Koyama & Inutsuka (2002),

$$\frac{\Lambda}{\Gamma} = 10^7 \exp\left(-\frac{118400}{T+1000}\right) + 1.4\times10^{-2}\sqrt{T}\exp\left(-\frac{92}{T}\right), \tag{4}$$

and $\Gamma = 2 \times 10^{-26}$ erg s^{-1}, which are taking into account many processes, e.g., photoelectric heating from dust grains and PAHs, heating by cosmic rays and X-rays, and atomic line cooling from hydrogen Lyα and C II (Koyama & Inutsuka 2000). Throughout this paper, we ignore the effect of magnetic fields. This might alter the results because the conductive heat flux is restricted along magnetic field lines, and therefore it should be taken into account in future analysis.

Before performing full numerical simulations, we show the results for a quasi-steady state (QSS), under which we transform the time derivative to the spatial one, $\partial/\partial t \rightarrow -\dot{R}_c\partial/\partial r$, where R_c is the cloud size and $\dot{R}_c$ is the velocity of the cloud size changing. Then the above equations become a set of ordinary differential equations. The details of the method to compute the structure will be found in a separate paper (Nagashima, Koyama & Inutsuka 2006). In Fig.1, several snapshots of evaporation rates, $\dot{M}(r) = 4\pi r^2\rho(r)v(r)$, are shown by thin dashed lines for the case of the saturation pressure, $p = p_{\rm sat} \simeq 2823k_B$. In contrast to the case of plane-parallel geometry, the cloud evaporates by the heat conduction and the size decreases as time passes. As we shall show in §4, the motion of the front shown in this figure is purely driven by the curvature effect (Nagashima, Koyama & Inutsuka 2005).

Next, we show the results obtained by full-numerical simulations. The code is based on the second-order Godunov method (van Leer 1977) in the Lagrange coordinate. We impose a constant pressure at the outer boundary, that is, the WNM pressure is given as the boundary condition. The initial structure of the cloud is computed by the QSS model with the size of about 0.1 pc. Correctly giving the initial condition is very important for this computation. If the deviation of the initial condition from the correct structure is large, spherical sound waves emerge and we could not extract the information on the evaporation. Several snapshots at which the size of the cloud becomes the same as that shown for the QSS are indicated by the thin solid lines in Fig.1. It is evident that the
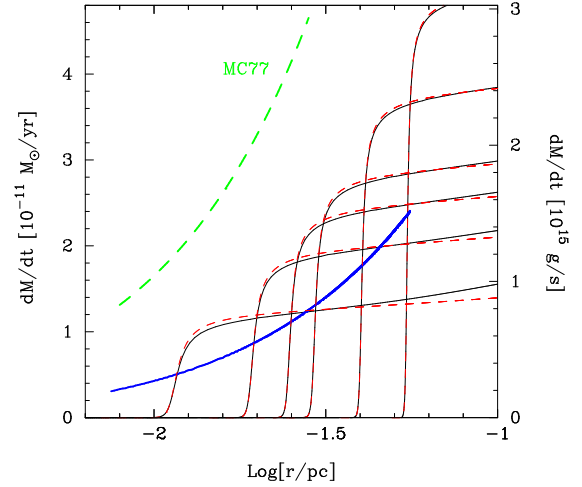

FIG. 1.— Evaporation rate, $\dot{M}$ for the case of the saturation pressure. Thin solid and dashed lines represent snapshots of $\dot{M} = 4\pi r^2\rho(r)v(r)$ as a function of the distance from the cloud center for the full numerical simulation and the QSS calculation, respectively. Thick solid curve indicates $\dot{M}$ as a function of the cloud size. The size of clouds are defined in the text. Thick dashed curve indicates $\dot{M}$ given by MC77.

results given by the full simulation agree well with those given by assuming the QSS probably because the Mach number of the flow is much below unity. Therefore we can consider that the QSS is a good assumption for the description of the evaporation of clouds.

The thick solid curve in Fig.1 denotes the evaporation rate as a function of the cloud size, $\dot{M}(R_c)$, where the edge of the cloud is defined as a radius at which the density becomes $(n_{\rm CNM} + n_{\rm WNM})/2$. We have confirmed that the rate is proportional to the size, $\dot{M}(R_c) \propto R_c$. In fact, this is the same scaling against the size as that given by MC77. Their evaporation rate is

$$\dot{M}_{\rm MC77} = \frac{16\pi\mu\kappa(T_f)R_c}{5k_B} = 1.3 \times 10^{15}T_f^{1/2}\left(\frac{R_c}{\rm pc}\right) {\rm g\,s^{-1}}, \tag{5}$$

where T_f is the temperature of the WNM, $T_f \simeq 6.4 \times 10^3$ K for the adopted heat-loss function. This rate is plotted by the thick dashed curve. When we write our evaporation rate in a similar way,

$$\dot{M} = 3.1 \times 10^{14}T_f^{1/2}\left(\frac{R_c}{\rm pc}\right) {\rm g\,s^{-1}}, \tag{6}$$

therefore this is about a factor of four smaller than MC77's in spite of the same dependence on the size. The form of the evaporation rate shall be discussed in §4 in a different way from MC77. Note that the above rate is valid only for the case of the saturation pressure. In the next section, we show the evaporation timescale for other cases.

3. EVAPORATION TIMESCALE

Hereafter we define the evaporation timescale as $t_{\rm evap} \equiv M/\dot{M}$, where the cloud mass M is estimated by integrating the mass density within the cloud radius, R_c, therefore $M_c \simeq 4\pi\rho_{\rm CNM}R_c^3/3$. Fig.2 shows the evaporation timescale against the cloud size R_c. The solid straight line indicates $t_{\rm evap}$ for the case of the saturation pressure, $p = p_{\rm sat}$, and for the QSS. The crosses on the solid line denote the result from the full numerical simulation. Similar to Fig.1, both are in excellent
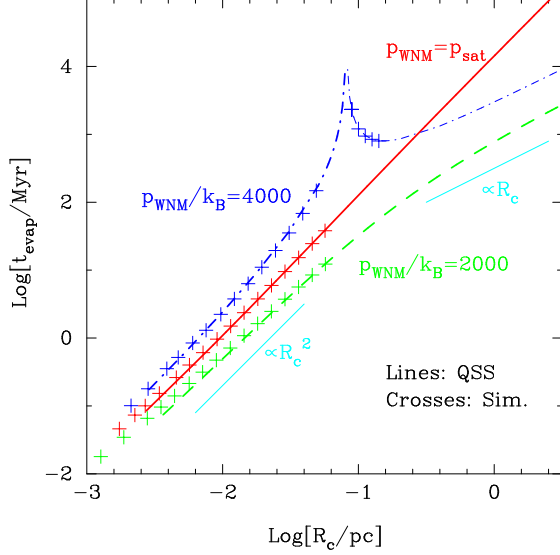
FIG. 2.— Evaporation timescales defined as $M_c/\dot{M}$ as a function of the radius of clouds. Solid, dashed, and dot-dashed lines indicate the timescale for $p = p_{\text{sat}}$, $2000k_B$, and $4000k_B$, respectively, under the assumption of the QSS. Crosses denote simulation results. Note that $R_c \gtrsim 0.1$ for $p_{\text{WNM}}/k_B = 4000$ corresponds to the condensation timescale.
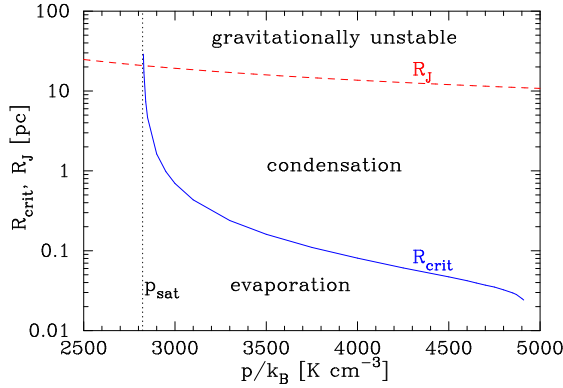

FIG. 3.— Critical radius R_{crit} plotted by the solid line against pressure of the WNM. Jeans length R_J is also plotted by the dashed line.

agreement. It is easily confirmed that $t_{\text{evap}} \propto R_c^2 \propto M^{2/3}$ for $p = p_{\text{sat}}$, reflecting the previous result $\dot{M} \propto R_c$ and $M \propto R_c^3$. Here it should be worth noting that tiny H I clouds with the size of $\sim 10^{-2}$ pc must disappear in ~ 1Myr irrespective of the ambient pressure.

The dashed and dot-dashed lines indicate the evaporation timescale for the cases of $p/k_B = 2000$ and 4000 K cm^{-3}, respectively. Crosses on the lines also denote the results from the full numerical simulations. Again we can see the good agreement. Note that clouds always evaporate for $p \leq p_{\text{sat}}$, but large clouds can grow for $p > p_{\text{sat}}$, which is shown by the thin dot-dashed line at $R \gtrsim 0.1$ pc. This means that there exists a critical radius for growth (condensation), R_{crit} (see the next section). Hereafter we use t_{evap} also as the growth timescale. For these pressures, the size dependence of t_{evap} is partly different from that for the saturation pressure. While it is very similar to that for the saturation pressure, $t_{\text{evap}} \propto R_c^2$ at $R_c \ll R_{\text{crit}}$, it becomes $t_{\text{evap}} \propto R_c$ at $R_c \gg R_{\text{crit}}$. This is very different feature from the expectation from MC77. In the next section, we explain what the size dependence comes from.

4. DISCUSSION

Below we derive an approximate evaporation rate in the same way as that of Nagashima, Koyama & Inutsuka (2005). Firstly, we assume that the evolution is isobaric because the motion of fluids is much slower than the sound speed. Then the energy equation (3) can be described as an evolution equation of enthalpy,

$$\frac{\gamma}{\gamma - 1}\frac{k_B}{\mu}\rho\left[\frac{\partial}{\partial t} + v\frac{\partial}{\partial r}\right]T = -\rho\mathcal{L} + \frac{1}{r^2}\frac{\partial}{\partial r}r^2\kappa\frac{\partial T}{\partial r}. \quad (7)$$

Next, again because of the slow motion of fluids, we assume that the fluids are in a QSS, so that the time derivative $\partial/\partial t$ can be replaced with $-\dot{R}_c\partial/\partial r$,

$$\frac{\gamma}{\gamma - 1}\frac{k_B}{\mu}\rho u\frac{\partial T}{\partial r} = -\rho\mathcal{L} + \frac{\partial}{\partial r}\kappa\frac{\partial T}{\partial r} + \frac{2}{r}\kappa\frac{\partial T}{\partial r}, \quad (8)$$

where $u \equiv v - \dot{R}_c$. The first and second terms in the right hand side emerge independent of geometry of fronts. Therefore, if it is reasonable to assume that the structure of fronts is almost independent of geometry, we can replace the two terms with the left hand side for the case of the plane-parallel geometry,

$$\frac{\gamma}{\gamma - 1}\frac{k_B}{\mu}\rho u\frac{\partial T}{\partial r} \simeq \frac{\gamma}{\gamma - 1}\frac{k_B}{\mu}\rho u_1(p)\frac{\partial T}{\partial r} + \frac{2}{r}\kappa\frac{\partial T}{\partial r}, \quad (9)$$

where $u_1(p)$ is the fluid velocity with respect to the front rest frame, and depends only on pressure. Noting the fact that the derivative of temperature, $\partial T/\partial r$ has a non-zero, finite value only at the front, we can omit the derivative and substitute R_c for r,

$$u(R) \simeq u_1(p) + \frac{\gamma - 1}{\gamma}\frac{\mu}{k_B}\frac{2}{R_c}\frac{\kappa_R}{\rho_R}, \quad (10)$$

where the subscript R stands for values at the front, $r = R_c$.

To obtain a formula for the size evolution, we need a further assumption that a convergence factor from u to $\dot{R}_c$ is independent of geometry, $f \equiv -\dot{R}_c/u$. This becomes exact for the case of the thin-front limit. For three-dimensional spherical geometry, from the mass-flux conservation, we obtain $-4\pi r^2\rho_{\text{CNM}}\dot{R}_c = 4\pi R_c^2\rho_R u_R$, where $v = 0$ inside the cloud. Then $f_{\text{sphere}} = (R_c^2\rho_R)/(r^2\rho_{\text{CNM}})$. For the plane-parallel geometry, $-\rho_{\text{CNM}}c = \rho_R u_R$, where c is the speed of the front with respect to the cloud (CNM) rest frame, which corresponds to $\dot{R}_c$. Then $f_{\text{p.p.}} = \rho_R/\rho_{\text{CNM}}$. If the front is sufficiently thin, and therefore the density ρ is ρ_{CNM} even just inside the front, we obtain $f_{\text{sphere}} \to \rho_R/\rho_{\text{CNM}}$. This is equal to $f_{\text{p.p.}}$. In reality, however, the front has a non-zero thickness. So the use of the same f is an approximation.

Finally, we thus obtain the size evolution,

$$\dot{R}_c \simeq c(p) - f\frac{\gamma - 1}{\gamma}\frac{\mu}{k_B}\frac{2}{R_c}\frac{\kappa_R}{\rho_R}, \quad (11)$$

where we explicitly write the argument of c, p, to stress that it depends only on pressure. For $p > p_{\text{sat}}$, $c(p) > 0$, and vice versa. This equation tells us the existence of a critical radius at which $\dot{R}_c = 0$ when $c(p) > 0$,

$$R_{\text{crit}} \simeq f\frac{\gamma - 1}{\gamma}\frac{\mu}{k_B}\frac{2}{c(p)}\frac{\kappa_R}{\rho_R}. \quad (12)$$

The evaporation rate is obtained from the above equation,

$$\dot{M} \simeq -4\pi R_c^2\rho_{\text{CNM}}\dot{R}_c \equiv \dot{M}_{\text{p}} + \dot{M}_c, \quad (13)$$

where $\dot{M}_\mathrm{p}$ and $\dot{M}_\mathrm{c}$ are the "pressure" and "curvature" terms, respectively, and

$$\dot{M}_\mathrm{p} = -4\pi R_c^2 \rho_\mathrm{CNM} c(p) \propto R_c^2, \qquad (14)$$

$$\dot{M}_\mathrm{c} = 4\pi R_c^2 \rho_\mathrm{CNM} f \frac{\gamma-1}{\gamma} \frac{\mu}{k_B} \frac{2}{R_c} \frac{\kappa_R}{\rho_{R_c}} \propto R_c. \qquad (15)$$

From the different dependences on R_c, we can see that the pressure term dominates for large clouds, $R_c \gg R_\mathrm{crit}$, and that the curvature term does for small clouds, $R_c \ll R_\mathrm{crit}$. Using $f \simeq \rho_R / \rho_\mathrm{CNM}$ and $\gamma = 5/3$, we obtain

$$\dot{M}_\mathrm{c} = \frac{16\pi\mu\kappa_R R_c}{5 k_B}. \qquad (16)$$

Thus the difference from MC77's result for the curvature term is

$$\frac{\dot{M}_c}{\dot{M}_\mathrm{MC77}} = \frac{\kappa_R}{\kappa(T_f)} = \sqrt{\frac{T_R}{T_f}}, \qquad (17)$$

because of $\kappa \propto \sqrt{T}$. Thus the difference between our result and MC77's can be partly explained by the above in addition to the thin-front approximation to derive the curvature term, $\dot{M}_\mathrm{c}$. MC77 took into account only the curvature term.

The size dependence of the evaporation timescale is thus dependent on the size. For small clouds, $R_c \ll R_\mathrm{crit}$, the timescale is $t_\mathrm{evap} \sim M/\dot{M}_\mathrm{c} \propto R_c^2$. For large clouds, $R_c \gg R_\mathrm{crit}$, it is $t_\mathrm{evap} \sim M/\dot{M}_\mathrm{p} \propto R_c$. Thus the size dependence of the evaporation timescale shown in Fig.2 is well explained by the above approximate evaporation rate we derived.

A further important point is the relationship between the critical radius, R_crit, and the Jeans length for gravitational instability, $R_\mathrm{J} \equiv c_S \sqrt{\pi/G\rho}$, where c_S is the sound velocity. Fig.3 shows R_crit and R_J as a function of the ambient pressure. We find that there can exist clouds growing only by condensation for the case of $p > p_\mathrm{sat}$. The timescale until the cloud size exceeds R_J is, however, quite long, as shown in Fig.2.

5. CONCLUSION

We have investigated the evaporation rate and timescale of tiny H I clouds by using numerical simulations. We confirmed that the results are almost recovered by assuming the evolution under the quasi-steady state because the fluid velocity is much slower than the sound velocity. We have found

that clouds with a size of about 0.01 pc evaporate in approximately 1Myr almost independent of pressure of the ambient WNM. This suggests that there might be some mechanisms to continuously form tiny H I clouds, or, that the ambient pressure around the clouds is much higher than the standard ISM pressure, if their existence is ubiquitous. The evaporation timescale of clouds larger than 0.1 pc, however, depends strongly on the pressure. Clouds larger than a critical radius can even grow if the pressure is higher than the saturation pressure.

In order to physically understand the simulation results, we derived an analytic formula for evaporation by assuming the isobaric evolution and a structure of the interface independent of geometry. The obtained evaporation rate has two separate terms: "pressure" and "curvature" terms. The former is independent of geometry, that is, it emerges even in analysis for plane-parallel geometry, and it becomes zero for the saturation pressure. The latter is proportional to the size of clouds, that is, the curvature. We have found that the obtained evaporation rate is a natural extension of an evaporation rate obtained by MC77, which corresponds to our "curvature" term, but MC77's rate is a little higher than ours. For lower pressure than the saturation pressure, the signs of the two term are the same. Therefore clouds always evaporate. For higher pressure, on the other hand, the signs are different. Therefore a critical size exists and the size of clouds determines whether they grow or evaporate. We have confirmed that the simulation results show the same dependence on the size as the analytic formula.

We would like to stress that our analysis can be adapted irrespective of the origin of the tiny H I clouds. In order to obtain information on the formation process, statistical properties such as mass function of clouds should be required. This will be done in future.

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